

Gravitational modulated reheating and non-Gaussianity in supergravity R^2 inflation

Yuki Watanabe¹ and Jun'ichi Yokoyama^{1,2}

¹*Research Center for the Early Universe (RESCEU),
Graduate School of Science, The University of Tokyo, Tokyo 113-0033, Japan*

²*Kavli Institute for the Physics and Mathematics of the Universe (WPI),
The University of Tokyo, Kashiwa, Chiba 277-8568, Japan*

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Abstract

Reheating after R^2 inflation proceeds through gravitational particle production of conformally noninvariant fields. We argue that the nonvanishing expectation value of flat directions generic in supersymmetric theories break conformal invariance of the fields coupled to them in a position-dependent manner due to quantum fluctuations. As a result modulated reheating can occur after the supergravity R^2 -inflation. The resultant curvature fluctuation is a mixture of the one produced during inflation and that produced by modulated reheating. The spectral index takes a value between $n_s = 0.960$ and 0.983 , the nonlinearity parameter of the local-type non-Gaussianity can be $f_{NL} \sim \pm 10$, and the tensor-to-scalar ratio is $r \leq 4 \times 10^{-3}$.

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1 Introduction

The R^2 inflation [1] is a unique inflation model in the sense that it is one of the oldest models of inflation [2] but still observationally viable. In fact, it is not only viable but preferred by the observational results of WMAP9 [3],¹ with its predictions of scalar spectral index,² $n_s = 0.964$, and the tensor-to-scalar ratio, $r = 3.9 \times 10^{-3}$, occupying the center of the likelihood contour. In order to further probe or falsify this model, it is desired to have more predictions that can be tested soon. Among such quantities is the nonlinearity parameter f_{NL} of curvature fluctuations [5] which measures the deviation of their statistical distribution from Gaussian.

Previously, R^2 inflation has been known to yield highly Gaussian curvature fluctuations, which can be easily understood regarding the extra scalar degree of freedom in $f(R)$ -type theory, dubbed *scalaron* φ , as the driving force of inflation [6]. Indeed φ has a very flat effective potential ideal for chaotic inflation [7] with a canonical kinetic term. In this paper we argue that nontrivial processes in the reheating phase after inflation may produce additional curvature perturbations, which can be appreciably non-Gaussian, through the so-called modulated reheating scenario [8] known to produce f_{NL} of local type [9, 10].

In this theory, inflaton is followed by damped oscillation of the scalar curvature $R(t)$, which induces gravitational particle production of conformally non-invariant fields thereby reheating the Universe. This process is equivalently described by the decay of the scalaron-inflaton φ which oscillates around the potential minimum after inflation. In the scalaron picture, φ is coupled with conformally non-invariant fields only, so it is coupled with neither massless fermions nor gauge bosons (except gauge conformal anomaly [11]). Therefore, the previous study of reheating after R^2 inflation [12, 13, 14] including more general $f(R)$ inflation [15] mostly considered creation of massless minimally coupled scalar bosons.

However, since mass terms break the conformal invariance, scalarons can decay into massive particles through their mass terms [16] (see also [17, 18] for earlier works). These channels open up a new possibility to produce curvature perturbations because the mass terms, which are determined by the Higgs field, turn out to be spatially dependent reflecting quantum fluctuations of the Higgs condensation [19] generated during inflation. As a result the modulated reheating [8] may take place which contributes to the observed curvature perturbations.

Unfortunately, however, such a mechanism does not work in the original R^2 inflation model which is realized by adding a single term, $R^2/(6M^2)$, to the gravitational Lagrangian. In this simplest model the scalaron mass is determined (or bounded from above, in case there is other sources of curvature perturbation besides the scalaron's fluctuations) by the observed amplitude of curvature perturbations as $M \simeq 3 \times 10^{13}$ GeV. As a result, its decay rate per mode $\Gamma_1 = M^3/(192\pi M_{\text{Pl}}^2) \simeq 8$ GeV is so small that the Universe would not be reheated until long after inflation.

¹The latest observational results of Planck [32, 33] also support the R^2 inflation, with tighter constraints on n_s , r and f_{NL} . We discuss their results further in subsection 4.1.

²The first prediction of the scalar spectral index for R^2 inflation was made correctly by Mukhanov and Chibisov [4]

During this long period of scalaron oscillation, the Higgs condensation also oscillates and decreases its amplitude. Therefore, by the time the reheating occurs, the standard model (SM) Higgs field value would be too small to leave any observational trace in contrast to the case studied in [22]. So much is the story of R^2 inflation in the SM.

On the other hand, it is barely likely that the SM is the ultimate theory with which we should describe the birth and evolution of the Universe. As the primary extension of the SM, there are a number of sound motivations to introduce the supersymmetry (SUSY). Then the R^2 inflation – or its analogue – should be realized in the context of local supersymmetry or the supergravity (SUGRA). An attempt to realize R^2 inflation in this context was first made by Ketov [20], and was improved by Ketov and Starobinsky [21]. Its most wonderful feature is that the scalaron mass in the reheating phase is not directly related to the amplitude of curvature fluctuations and can be much heavier than in the original model, so that the Universe can be reheated immediately after inflation by gravitational particle production. The other plausible feature is the existence of flat-direction scalar fields which have potentials much flatter than the SM Higgs field has, even after the effects of SUSY breaking and non-renormalizable superpotential are taken into account. Thus quantum fluctuations around a large expectation value of the flat direction can induce modulated reheating [23], and generation of appreciable amount of local-type non-Gaussianity is possible in this model.

The rest of the paper is organized as follows. In Sec. 2 we introduce R^2 inflation in SUGRA proposed in [21] and further analyzed in [24]. Then in Sec. 3 we incorporate a SUSY flat direction ϕ to analyze non-Gaussianity produced in this model. Sec. 4 is devoted to conclusion. In the appendix we describe scalaron interactions with other fields and calculate its decay widths as an equivalent way to calculate gravitational particle production rate.

2 R^2 inflation in SUGRA

Here first we briefly introduce higher curvature inflation model proposed by Ketov [20] and Ketov and Starobinsky [21] in SUGRA. The theory is defined by the action

$$S = \kappa^{-1} \int d^4x d^2\theta \mathcal{E} F(\mathcal{R}) + \text{h.c.}, \quad (1)$$

where $F(\mathcal{R})$ is a function of the scalar curvature superfield

$$\begin{aligned} \mathcal{R} = & -\frac{\kappa}{3} B^* - \theta \left(\sigma^a \bar{\sigma}^b \psi_{ab} - i \sigma^a \bar{\psi}_a \frac{\kappa}{3} B^* + i \psi_a b^a \right) \\ & + \theta \theta \left(\frac{\kappa}{3} R - \frac{i\kappa}{2} \epsilon^{abcd} R_{abcd} - \frac{4\kappa^3}{9} B B^* + \dots \right), \end{aligned} \quad (2)$$

and \mathcal{E} is the chiral superspace density in a Wess-Zumino type gauge

$$\mathcal{E} = e(x) \left[1 - 2i\theta \sigma_a \bar{\psi}^a(x) + \theta \theta \kappa^2 B(x) \right]. \quad (3)$$

Here, $e = \sqrt{-g}$, ψ^a is the gravitino, B is an auxiliary scalar field, and κ is the reciprocal of the reduced Planck scale $M_{Pl} = \kappa^{-1} = m_{Pl}/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV.

Ignoring fermions, we find

$$F(\mathcal{R}) = F\left(-\frac{\kappa}{3}B^*\right) + F'\left(-\frac{\kappa}{3}B^*\right)\left(\frac{\kappa}{3}R - \frac{4\kappa^3}{9}BB^*\right)\theta\theta, \quad (4)$$

$$S = \int d^4x e \left[-3XF(X^*) + F'(X^*)\left(\frac{1}{3}R - 4XX^*\right) \right] + \text{h.c.}, \quad (5)$$

$$X \equiv -\frac{\kappa}{3}B. \quad (6)$$

Ignoring a pseudo-scalar partner (axion) of the scalaron, namely, taking $X = X^*$ and $F = F^*$, the Lagrangian of the gravity sector reads,

$$\mathcal{L} = -6XF(X) + 2\left(\frac{1}{3}R - 4X^2\right)F'(X), \quad (7)$$

together with the constraint which is nothing but the equation of motion (EOM) for the auxiliary field

$$0 = 3F(X) + 11XF'(X) - \left(\frac{1}{3}R - 4X^2\right)F''(X). \quad (8)$$

The choice of Ketov and Starobinsky is

$$F(\mathcal{R}) = \frac{1}{2}f_1\mathcal{R} + \frac{1}{2}f_2\mathcal{R}^2 + \frac{1}{6}f_3\mathcal{R}^3, \quad (9)$$

which leads the Lagrangian

$$\mathcal{L} = \frac{1}{3}f_1R + \frac{2}{3}f_2RX + \left(\frac{1}{3}f_3R - 7f_1\right)X^2 - 11f_2X^3 - 5f_3X^4, \quad (10)$$

and the constraint

$$0 = X^3 + \frac{33f_2}{20f_3}X^2 - \frac{1}{30}(R - R_0)X - \frac{f_2}{30f_3}R, \quad (11)$$

with

$$R_0 \equiv \frac{21f_1}{f_3}, \quad (12)$$

where f_1 , f_2 , and f_3 are positive constants. Note that we use the metric signature $(-, +, +, +)$ so that our sign convention of \mathcal{R} and $R = 12H^2 + 6\dot{H}$ is different from that of Ketov and Starobinsky.

Since the gravitational constant must be positive for stability of the system, $F'(X) > 0$ and thus $f_1f_3 > f_2^2$. The absence of ghost and tachyonic degrees of freedom requires $f_1 > 0$ and $f_3 > 0$. The sub-Planckian curvature during inflation

requires $f_3 \gg 1$. The sub-Planckian scalaron mass after inflation requires $f_2^2 \gg f_1$ in order to avoid large quantum gravity loop corrections. In summary, [21]

$$f_3 \gg 1, \quad f_1 f_3 > f_2^2 \gg f_1 > 0. \quad (13)$$

In terms of the reduced Planck scale M_{Pl} , the scalaron mass during inflation, M , and that in the reheating stage around the potential minimum, m , they can be expressed as

$$f_1 = \frac{3}{2} M_{Pl}^2, \quad f_2 = \sqrt{\frac{63}{8}} \frac{M_{Pl}^2}{m}, \quad \text{and} \quad f_3 = \frac{15 M_{Pl}^2}{M^2}, \quad (14)$$

respectively. Then the conditions (13) read

$$m > \sqrt{\frac{7}{20}} M, \quad M \ll M_{Pl}, \quad m \ll M_{Pl}, \quad (15)$$

respectively [24].

In the high curvature regime $R \gg R_0 = 21f_1/f_3 = 21M^2/10$, the Lagrangian reads

$$\begin{aligned} \mathcal{L} &= \frac{1}{3} f_1 R + \frac{1}{180} f_3 R^2 + \frac{\sqrt{30}}{100} f_2 R^{3/2} \\ &= \frac{M_{Pl}^2}{2} \left(R + \frac{R^2}{6M^2} + \frac{3\sqrt{105}}{100} \frac{R^{3/2}}{m} \right). \end{aligned} \quad (16)$$

The dynamics of inflation realized in this Lagrangian has been studied by Ketov and Tsujikawa [24] in detail. They solved the EOM for the Hubble parameter of the Friedmann Universe, $H(t)$, and expressed the number of e -folds, N , acquired during inflation after the epoch $H(t) = H$ in terms of the dimensionless quantities

$$\alpha \equiv \frac{M^2}{mH} \quad \text{and} \quad \beta \equiv \frac{M^2}{H^2}, \quad (17)$$

as

$$\begin{aligned} N &= \frac{1}{126\alpha^2} \left[3\alpha \left(80\sqrt{35} - 21\alpha - \sqrt{7(63\alpha^2 + 16000\beta/3)} \right) - \frac{4000\beta}{3} (8 \ln 2 + 3 \ln 5) \right. \\ &\quad \left. + \frac{8000\beta}{3} \ln \left(\frac{\sqrt{7(63\alpha^2 + 800\beta)} + 21\alpha\sqrt{63\alpha^2 + 1600\beta}}{21\alpha + 2\sqrt{35}\beta} \right) \right]. \end{aligned} \quad (18)$$

The square amplitude of curvature perturbation is given by

$$\mathcal{P}_{\zeta_{R2}}(N) \cong \frac{1250}{3\pi^2} \left(\frac{M}{M_{Pl}} \right)^2 \left(3\sqrt{35}\alpha + \frac{100\beta}{3} \right)^{-2}. \quad (19)$$

We assume this curvature perturbation generated by inflation is responsible for the fraction λ^2 of the total amplitude, $\mathcal{P}_{obs} \cong 2.4 \times 10^{-9}$, determined by COBE and WMAP [3] on the pivot scale $k_* = 0.002 \text{Mpc}^{-1}$, namely,

$$\mathcal{P}_{\zeta_{R2}}(N_*) = \lambda^2 \mathcal{P}_{obs} = 2.4 \times 10^{-9} \lambda^2, \quad (20)$$

where N_* is the number of e -folds of inflation after the comoving pivot scale left the Hubble radius.

The scalar spectral index n_s and the tensor-to-scalar ratio r can also be expressed by α and β as

$$n_s = 1 - \frac{3\sqrt{35}}{100}\alpha - \frac{2\beta}{3}, \quad \text{and} \quad r = \frac{1}{2500}(9\sqrt{35}\alpha + 100\beta)^2, \quad (21)$$

for the case $\lambda^2 = 1$ [24].

As is seen in (16), the original R^2 -inflation is obtained in the small α limit, when N approaches $3/\beta - 1/2$ or $\beta = 6/(2N + 1)$. Then the COBE-WMAP normalization yields

$$M = \frac{7.54 \times 10^{-4}}{N_* + 1/2} \lambda M_{Pl} = 1.36 \times 10^{-5} \lambda \nu_*^{-1} M_{Pl}, \quad \nu_* \equiv \frac{N_* + 1/2}{55.5} \quad (22)$$

which in turn determines the Hubble parameter as a function of N as

$$H = \left(\frac{N + 1/2}{6} \right)^{1/2} M. \quad (23)$$

Thus the Hubble parameter when the pivot scale left the Hubble radius is given by

$$H_* = 4.1 \times 10^{-5} \lambda \nu_*^{-1/2} M_{Pl}. \quad (24)$$

On the other hand, the mass of the scalaron in the reheating phase is determined as

$$m = \frac{M^2}{\alpha_* H_*} = 4.47 \times 10^{-6} \lambda \nu_*^{-3/2} \alpha_*^{-1} M_{Pl}, \quad (25)$$

so that the scalaron mass in the reheating phase is unrelated with the CMB observation, and can be chosen arbitrarily under the condition $m \ll M_{Pl}$. Here α_* is the value of α at $H = H_*$. Thus the decay rate of the scalaron (see Appendix A and set $M_s = m$ there),

$$\Gamma \cong \mathcal{N}_S \Gamma_{SK} = \frac{\mathcal{N}_S m^3}{192\pi M_{Pl}^2}, \quad (26)$$

where \mathcal{N}_S stands for the number of scalar decay modes due to scalar kinetic interaction terms, can be larger than the cosmic expansion rate at the end of inflation, $H_f \cong M/\sqrt{12}$, if

$$m > 0.133 \mathcal{N}_S^{-1/3} \nu_*^{-1/3} \lambda^{1/3} M_{Pl} \equiv m_{rh}. \quad (27)$$

In such cases, the Universe is reheated efficiently soon after inflation even if the inflaton-scalaron decays only through gravitational interaction. In what follows we fix $m = m_{rh}$ so that the Universe is reheated rapidly after inflation.

The above inequality (27) imposes an upper bound on α as

$$\alpha_* < 1.6 \times 10^{-4} \lambda^{2/3} \nu_*^{-7/6} (\mathcal{N}_S/10^2)^{1/3}. \quad (28)$$

Thus the rapid reheating after inflation automatically guarantees the original R^2 -inflation in the supergravity context. The scalar spectral index n_{sR2} and the tensor-to-scalar ratio r_{R2} from inflation are given by

$$n_{sR2} = 1 - \frac{4}{2N+1} = 0.964, \quad r_{R2} = \frac{48}{(2N+1)^2} = 3.9 \times 10^{-3}, \quad (29)$$

respectively, where we have quoted the numerical values at $N = 55$.

Finally we comment on the thermal history after inflation in this model. It has been known that finite amount of radiation is created after the onset of field oscillation even if only perturbative decay with the rate (26) operates, and its energy density is given by $\rho_r(t) = \frac{6}{5}\Gamma H(t)M_{Pl}^2$ (see e.g. [27]). Since Γ is as high as H_f here by assumption, the radiation temperature just after inflation is comparable—in fact only 20% lower—to the case all the energy density is immediately converted to radiation at the end of inflation. Thus even if the preheating may occur just after inflation as analyzed by Ketov and Tsujikawa [24], the thermal history does not change significantly. Whether preheating occurs or not in the presence of large Γ deserves further study, though. The high temperature effect may modify the potential of the flat direction to initiate field oscillation due to thermal effects, which does not change $\delta\phi/\phi$ because both the homogeneous mode ϕ and superhorizon long-wave fluctuations $\delta\phi$ would oscillate in the same manner. Furthermore, in our model, the scalaron decays rapidly before these oscillations dissipate field amplitude appreciably. So we may use the field values at the end of inflation below.

Since the reheat temperature is quite high, the gravitino problem is a problem [25]. We should either adopt a model with large enough gravitino mass without changing cosmic expansion history, or invoke non-standard cosmology like thermal inflation [26]. In the former case, the pivot scale is pushed up to $N_* = 60$ due to the earlier radiation domination, while in the latter case it may be shifted to a value as small as $N_* \sim 50$.

3 Modulated reheating through SUSY flat directions

We now consider behaviors of a SUSY flat direction field in this inflation model. Although our favorite one is $H_u H_d$ flat direction in close analogue to the Higgs condensation in the SM, other combinations such as LH_u or those involving three fields would also work with their appropriate choice of coupling constants, provided they do not suffer from Q-ball instability before the scalaron decay.

Here we denote a generic flat direction by ϕ and study its cosmological effects.³

³Note that a flat direction ϕ is not the inflaton, but causes curvature perturbations by modulated reheating. The $H_u H_d$ flat direction can be associated with it by ${}^t H_u \propto (0, \phi)$ and ${}^t H_d \propto (\phi, 0)$.

By definition ϕ acquires a potential only through SUSY breaking and possible non-renormalizable terms in the superpotential, $W \supset \frac{(\sqrt{2}\Phi)^n}{nM_X^{n-3}}$ with $\Phi = \frac{\phi}{\sqrt{2}}e^{i\delta}$. Since we are interested only in the (position-dependent) amplitude of the flat direction, we suppress possible CP violating A-terms to write the potential as

$$V[\phi] = \frac{1}{2}m_0^2\phi^2 + \frac{\phi^{2(n-1)}}{M_X^{2(n-3)}}, \quad n \geq 4, \quad (30)$$

for $\phi \ll M_{Pl}$. Here m_0 is the SUSY breaking mass assumed to be $\mathcal{O}(1\text{TeV})$ or so, and M_X is a cutoff scale to characterize nonrenormalizable interactions.

The second term is dominant for $\phi > m_0^{\frac{1}{n-2}} M_X^{\frac{n-3}{n-2}} \equiv \phi_{NR}$ which we assume is satisfied during inflation in the subsequent analysis. For $n = 4$ and 6 we find $\phi_{NR} = 2.0 \times 10^{-8} M_{Pl}$ and $1.4 \times 10^{-4} M_{Pl}$, respectively, for $m_0 = 1\text{TeV}$ and $M_X = M_{Pl}$. With no Hubble induced mass terms here during inflation, the initial value of ϕ is randomly distributed, and it does not evolve significantly during inflation if its effective mass is smaller than the Hubble parameter, $V''[\phi] < H_f^2$. This yields an upper bound on ϕ as

$$\phi < \left[\frac{H_f^2 M_X^{2n-6}}{(2n-2)(2n-3)} \right]^{\frac{1}{2n-4}} \equiv \phi_{\max}. \quad (31)$$

If it was not satisfied in the beginning, the field would roll down the potential until it is satisfied. Using (22) and (23), we find

$$\begin{aligned} \phi_{\max} &= 8.5 \times 10^{-4} \lambda^{1/2} \nu_*^{-1/2} \left(\frac{M_X}{M_{Pl}} \right)^{1/2} M_{Pl}, \quad (n=4), \\ \phi_{\max} &= 2.5 \times 10^{-2} \lambda^{1/4} \nu_*^{-1/4} \left(\frac{M_X}{M_{Pl}} \right)^{3/4} M_{Pl}, \quad (n=6), \end{aligned} \quad (32)$$

respectively.

It is convenient to express solutions of the homogeneous mode, ϕ , and fluctuation around it, $\delta\phi$, as a function of the number of e -folds until the end of inflation, N , which is a decreasing function of cosmic time, using the Hubble parameter (23) [10]. The slow-roll EOM of ϕ reads

$$\frac{d\phi}{dN} = \frac{1}{3 - H'/H} \frac{V_\phi}{H^2} = \frac{6V_\phi}{(3N+1)M^2}, \quad H' = \frac{dH}{dN}, \quad (33)$$

which is solved as

$$\phi(N) = \phi_* \left[1 + \frac{8(n-1)(n-2)\phi_*^{2n-4}}{M^2 M_X^{2n-6}} \ln \left(\frac{N_* + 1/3}{N + 1/3} \right) \right]^{\frac{-1}{2n-4}}. \quad (34)$$

Here ϕ_* is the initial value at $N = N_*$, which we take the epoch of the wave number of our interest, namely, the pivot scale k_* left the Hubble radius during inflation,

and a prime denotes differentiation with respect to N . In particular, at the end of inflation corresponding to $N = 0$, we find

$$\begin{aligned}\phi(0) &= \phi_* \left[1 + \frac{8(n-1)(n-2)\phi_*^{2n-4}}{M^2 M_X^{2n-6}} \ln(3N_* + 1) \right]^{\frac{-1}{2n-4}} \equiv \phi_f = \phi_*(1 + \Delta)^{\frac{-1}{2n-4}}, \\ \Delta &\equiv \frac{8(n-1)(n-2)\phi_*^{2n-4}}{M^2 M_X^{2n-6}} \ln(3N_* + 1) < \frac{n-2}{3(2n-3)} \ln(3N_* + 1) \equiv \Delta_{\max},\end{aligned}\quad (35)$$

where we have used (31) and $H_f \cong M/\sqrt{12}$. We find $\Delta_{\max}^{n=4} = 0.682$ and $\Delta_{\max}^{n=6} = 0.757$ for $N_* = 55$, respectively.

Similarly, the super-horizon evolution of fluctuation can be obtained by perturbatively solving the EOM

$$\frac{d\delta\phi}{dN} = \frac{1}{3 - H'/H} \left[\frac{V_{\phi\phi}}{H^2} \delta\phi + \frac{V_{\phi\phi\phi}}{2H^2} (\delta\phi)^2 \right], \quad (36)$$

to yield

$$\begin{aligned}\delta\phi(N) &= \delta\phi_* \exp \left(\int_{N_*}^N \frac{1}{3 - H'/H} \frac{V_{\phi\phi}}{H^2} dN \right) \\ &+ \frac{1}{2} (\delta\phi_*)^2 \int_{N_*}^N \frac{dN'}{3 - H'/H} \frac{V_{\phi\phi\phi}}{H^2} \exp \left(\int_{N_*}^{N'} \frac{1}{3 - H'/H} \frac{V_{\phi\phi}}{H^2} dN'' \right) \\ &\times \exp \left(\int_{N_*}^N \frac{1}{3 - H'/H} \frac{V_{\phi\phi}}{H^2} dN \right).\end{aligned}\quad (37)$$

Here the initial condition is given by $\delta\phi_* = H_*/(2\pi)$. At the end of inflation, we find

$$\delta\phi_f = \delta\phi_* \left(\frac{\phi_f}{\phi_*} \right)^{2n-3} + \frac{1}{2} \Theta (\delta\phi_*)^2, \quad (38)$$

where Θ is given by

$$\begin{aligned}\Theta &= -\frac{2n-3}{\phi_*^{2n-3}} \left(\phi_*^{2n-4} - \phi_f^{2n-4} \right) \left(\frac{\phi_f}{\phi_*} \right)^{2n-3} \\ &= -\frac{(2n-3)\Delta}{(1+\Delta)^{2+\frac{1}{n-2}} \phi_f} \equiv -\frac{\Upsilon}{\phi_f}.\end{aligned}\quad (39)$$

To calculate the power spectrum of curvature perturbation generated by modulated reheating, only the first term in (38) is important, whereas the second term is also necessary in principle to calculate non-Gaussianity.

Since the Universe is reheated rapidly after inflation, we can estimate the amplitude of curvature perturbation produced by the modulated reheating, ζ_m , and its non-Gaussianity using ϕ_f and $\delta\phi_f$. For this purpose it is convenient to use the scalaron picture and use the formulas given in [10].

First ζ_m is given by

$$\zeta_m = xQ'(x)\frac{\delta\Gamma}{\Gamma}, \quad (40)$$

where $Q(x)$ is defined by

$$Q(x) = Q(\Gamma/H) = \frac{1}{4} \ln \left[\int_0^\infty dN' \frac{\Gamma}{H(N')} e^{4N'} \frac{\rho_\varphi(N')}{\rho_{\varphi f}} \right]. \quad (41)$$

Here ρ_φ is the effective energy density of the scalaron-inflaton and the subscript f denotes the value at the end of inflation [10]. Our choice of the mass parameter corresponds to $x = 1$, where we find $Q'(1) = -0.0890$ and $Q''(1) = 0.126$. Here the total decay rate Γ is given by $\Gamma \cong \mathcal{N}_S \Gamma_{SK}$ ($\mathcal{N}_S = \mathcal{N}_\sigma + \mathcal{N}_\phi +$ all other light scalars) and the space-dependent part of the rate is expressed as

$$\begin{aligned} \delta\Gamma &= \Gamma_\phi \delta\phi = \mathcal{N}_\sigma \Gamma_{SKM} \phi \delta\phi + \mathcal{N}_\psi \Gamma_{FM} \phi \delta\phi \\ &= \frac{(\mathcal{N}_\sigma + \mathcal{N}_\psi) y^2 m \phi \delta\phi}{24\pi M_{Pl}^2} \equiv \frac{\mathcal{N}_{FL} y^2 m \phi \delta\phi}{24\pi M_{Pl}^2}, \end{aligned} \quad (42)$$

where we have put $h = y$ and replaced M_s by m in the formulas (76) and (77). Thus ζ_m is given by

$$\begin{aligned} -\zeta_m &= 0.712 \frac{\mathcal{N}_{FL} y^2 \phi_f \delta\phi_f}{\mathcal{N}_S^{1/3} m^2} = 2.33 \times 10^3 \mu \nu_*^{2/3} \lambda^{-2/3} \frac{\phi_f \delta\phi_f}{M_{Pl}^2}, \\ \mu &\equiv \left(\frac{y}{3}\right)^2 \left(\frac{\mathcal{N}_{FL}}{30}\right) \left(\frac{\mathcal{N}_S}{10^2}\right)^{-1/3}. \end{aligned} \quad (43)$$

Here μ is a numerical factor normalized by typical values of the number of decay modes and Yukawa coupling. Note that the number of decay modes for a pair of a single species of quark and antiquark, say t and \bar{t} , is $\mathcal{N}_\psi = 12$ and supersymmetry doubles it by the same contribution of Γ_{SKM} . We also note that Yukawa coupling is typically large, which might even saturate perturbative bound $y^2/(4\pi) \sim 1$ at high energy scale of our concern [28]. Using (35), (38) and (24), we find

$$-\zeta_m = 1.53 \times 10^{-2} \mu \nu_*^{1/6} \lambda^{1/3} (1 + \Delta)^{-\frac{n-1}{n-2}} \frac{\phi_*}{M_{Pl}}. \quad (44)$$

It should contribute to the fraction $1 - \lambda^2$ of the observed power spectrum by assumption, so $\mathcal{P}_{\zeta_m} = (1 - \lambda^2) \mathcal{P}_{obs}$. Then ϕ_*/M_{Pl} can be expressed as

$$\frac{\phi_*}{M_{Pl}} = 3.2 \times 10^{-3} \mu^{-1} \nu_*^{-1/6} \lambda^{-1/3} (1 - \lambda^2)^{1/2} (1 + \Delta)^{\frac{n-1}{n-2}}. \quad (45)$$

The spectral index, n_{sm} , of perturbation produced by modulated reheating is given by

$$n_{sm} \cong 1 - \frac{2}{2N + 1} = 0.982, \quad (46)$$

where the last equality is for the case $N = 55$.

Next we turn to the non-Gaussianity. It is known that modulated reheating produces an appreciable amount of non-Gaussianity of local type [9, 10], and the nonlinearity parameter f_{NL} of ζ_m , f_{NLm} , is given by

$$f_{NLm} = \frac{5Q''(x)}{6Q'^2(x)} + \frac{5}{6xQ'(x)} \left(\frac{\Gamma\Gamma_{\phi\phi}}{\Gamma_\phi^2} + \frac{\Gamma\Theta}{\Gamma_\phi} \right) = 13.2 - 1.17(1 - \Upsilon) \frac{\mathcal{N}_S m^2}{\mathcal{N}_{FL} y^2 \phi_f^2} \quad (47)$$

for $x = 1$. The full nonlinearity parameter is given by

$$\begin{aligned} f_{NL} &= f_{NLm} \frac{\mathcal{P}_{\zeta_m}^2}{\mathcal{P}_{obs}^2} = f_{NLm} (1 - \lambda^2)^2 \\ &= 13.2(1 - \lambda^2)^2 + 35.0(\Upsilon - 1)(1 + \Delta)^{-\frac{2n-3}{n-2}} \mu \nu_*^{-1/3} \lambda^{4/3} (1 - \lambda^2), \end{aligned} \quad (48)$$

where we have used (27), (35) and (44).

From (35), (45) and (22) Δ satisfies

$$\Delta = 1.38 \times 10^2 \frac{\ln(3N_* + 1)}{\ln 166} \nu_*^{4/3} \mu^{-4} (1 + \Delta)^6 \lambda^{-10/3} (1 - \lambda^2)^2 \left(\frac{M_{Pl}}{M_X} \right)^2 \quad (49)$$

for $n = 4$, and

$$\Delta = 4.78 \times 10^{-8} \frac{\ln(3N_* + 1)}{\ln 166} \nu_*^{2/3} \mu^{-8} (1 + \Delta)^{10} \lambda^{-14/3} (1 - \lambda^2)^4 \left(\frac{M_{Pl}}{M_X} \right)^6 \quad (50)$$

for $n = 6$. Here note that we are not allowed to take $\lambda = 0$ exactly, because we have first fixed M from (22) and this procedure is invalid there. Since this problem arises if and only if we take λ exactly equal to zero from the beginning, our formula for f_{NL} and n_s below are applicable even in the limit λ is small and approaches to zero.

For definiteness, let us take $\nu_* = 1$ (i.e. $N_* = 55$) and $\mu = 1$ below. Then we can express f_{NL} as a function of λ and Δ for each n as

$$f_{NL} = 13.2(1 - \lambda^2)^2 + 35.0K(n, \Delta) \lambda^{4/3} (1 - \lambda^2), \quad (51)$$

where

$$K(n, \Delta) \equiv (\Upsilon - 1)(1 + \Delta)^{-\frac{2n-3}{n-2}} = (2n - 3)\Delta(1 + \Delta)^{-\frac{4n-6}{n-2}} - (1 + \Delta)^{-\frac{2n-3}{n-2}}. \quad (52)$$

We can also determine the cut off scale from (49) and (50) as

$$M_X = 12\Delta^{-1/2}(1 + \Delta)^3 \lambda^{-5/3} (1 - \lambda^2) M_{Pl}, \quad (n = 4), \quad (53)$$

$$M_X = 6.0 \times 10^{-2} \Delta^{-1/6} (1 + \Delta)^{5/3} \lambda^{-7/9} (1 - \lambda^2)^{2/3} M_{Pl}, \quad (n = 6), \quad (54)$$

respectively.

Now let us evaluate $K(n, \Delta)$ for $n = 4$ and $n = 6$. We find

$$-1 \leq K(4, \Delta) \leq -1.93 \times 10^{-2}, \quad (55)$$

which is monotonically increasing as Δ , and the inequalities are saturated at $\Delta = 0$ and $\Delta_{\text{max}} = 0.682$, respectively. For $n = 6$ we find

$$-1 \leq K(6, \Delta) \leq 0.327, \quad (56)$$

where the minimum is realized again at $\Delta = 0$ and the maximum at $\Delta = 0.451$.

The case $\Delta \rightarrow 0$ corresponds to the limit M_X is large and nonrenormalizable interaction is absent. Then, independent of n we find

$$f_{NL} = 13.2(1 - \lambda^2)^2 - 35.0\lambda^{4/3}(1 - \lambda^2), \quad (57)$$

which takes maximum $f_{NL} = 13.2$ in the limit $\lambda \rightarrow 0$, and minimum $f_{NL} = -7.9$ at $\lambda = 0.749$. The former corresponds to the case all the observed fluctuations are due to the modulated reheating with the spectral index $n_s = 0.982$, while at $\lambda = 0.749$ we find $n_s = 0.972$ for $\nu_* = 1$, because the spectral index of total curvature perturbation, n_s is given by

$$n_s = \lambda^2 n_{sR2} + (1 - \lambda^2) n_{sm} = 1 - \frac{2\lambda^2 + 2}{111\nu_*}, \quad (58)$$

in general.

On the other hand, with $K(4, 0.682) = -1.93 \times 10^{-2}$, one finds

$$f_{NL} = 13.2(1 - \lambda^2)^2 - 0.675\lambda^{4/3}(1 - \lambda^2). \quad (59)$$

In this case f_{NL} decreases monotonically from 13.2 to -8.32×10^{-3} as we increase λ from 0 to 0.988. Then f_{NL} increases to 0 as $\lambda \rightarrow 1$.

For the case $K(6, 0.451) = 0.327$ we have

$$f_{NL} = 13.2(1 - \lambda^2)^2 + 11.5\lambda^{4/3}(1 - \lambda^2), \quad (60)$$

which takes maximum $f_{NL} = 13.5$ at $\lambda = 0.148$ and then decreases to 0 as λ is increased.

Our results suggest that as far as f_{NL} parameter is concerned, the nonrenormalizable term in the potential of the flat direction does not have observationally significant effects, which means that our result is robust in this respect.

If we take smaller values for μ , then the second term in (48) contributes less and the overall f_{NL} can become larger for $n = 4$ and smaller for $n = 6$ than the demonstrated case above.

4 Conclusion

In the present paper we have reconsidered cosmic history after R^2 -inflation in SUGRA in which reheating proceeds through gravitational particle production of conformally noninvariant fields. We have argued that conformal invariance is broken through a nonvanishing expectation value of a SUSY flat direction which gives a position dependent mass for those fields coupled to it due to its long-wave quantum fluctuations acquired during inflation. As a result such fluctuations induce modulated reheating.

We have shown that with reasonable values of the model parameters, on condition that the universe is reheated rapidly, curvature perturbation produced by modulated reheating may contribute to the observed fluctuation significantly and yields a local non-Gaussianity with $f_{NL} \sim \pm 10$.

The intriguing feature of our model is that the observable parameters n_s , r , f_{NL} are mutually correlated with each other through the ratio of the two contributions λ^2 , although the exact one-to-one correspondence is difficult due to the uncertainties in the pivot scale ν_* and the particle-physics parameter μ . For example, from (58) and

$$r = r_{R^2} \lambda^2 = \frac{48\lambda^2}{(2N_* + 1)^2} = \frac{48\lambda^2}{111^2 \nu_*^2}, \quad (61)$$

one can obtain λ and ν_* (or N_*), which can be compared with f_{NL} . In future, information on the thermal history carried by ν_* can also be tested by space based laser interferometer such as DECIGO [29, 30], and can be compared with our result.

4.1 Note added after Planck 2013 results

After we submitted the original version of this article to the arXiv, the Planck collaboration released their data. It gives [32, 33]

$$n_s = 0.9653 \pm 0.0069, \quad r < 0.13 \quad (95\% \text{ CL}), \quad f_{NL} = 2.7 \pm 5.8 \quad (68\% \text{ CL}), \quad (62)$$

where we have quoted values from Planck combined with WMAP polarization and lensing data sets for n_s and r . Their results are basically in good agreement with the predictions of our model.

To be precise, the lower value of n_s , which favors larger λ , does not yield any sizable non-Gaussianity with our model. On the other hand, the higher allowed value like $n_s = 0.967$ yields

$$r = 3.2 \times 10^{-3}, \quad f_{NL} = \{-4.8, 0.28, 2.1\} \quad (63)$$

for $\{K(n, 0_+), K(4, 0.68), K(6, 0.45)\}$, $\lambda = 0.91$, $\mu = 1$ and $\nu_* = 1$. Thus our model covers lower allowed range of the observed f_{NL} , namely, between -0.1σ and -1.3σ off the mean value.

From (58) one can see that smaller ν_* (or N_*) allows smaller λ resulting in higher values for $|f_{NL}|$. The thermal inflation scenario [26] after R^2 inflation may yield such values.

Finally, let us comment on the SM Higgs inflation with nonminimal coupling to gravity [34], which is also preferred by the Planck data [32]. As was emphasized in [35], predictions of n_s and r from the Higgs inflation and R^2 inflation are almost degenerate. Only the slight difference comes from the reheating stage. In fact, the SUSY flat direction alters reheating significantly from that after the original R^2 inflation, thereby producing non-negligible f_{NL} as opposed to the Higgs inflation. In our scenario the reheating temperature can be as high as 10^{14} GeV and spatially modulated.

A Scalon interaction with matter fields

In generic $f(R)$ -type gravitational theory,

$$S = \frac{1}{2\kappa^2} \int f(R) \sqrt{-g} d^4x + S_m, \quad (64)$$

the extra scalar degree of freedom, φ , of gravity is extracted by

$$\varphi \equiv \sqrt{\frac{3}{2}} M_{Pl} \ln |f'(R)|, \quad (65)$$

and its mass M_s is given by $M_s = 1/\sqrt{3f''(R)}$. So in the original R^2 -inflation model with $f(R) = R + R^2/(6M^2)$, the scalaron mass is constant and given by M which is fixed by the observed amplitude of curvature fluctuations. The supergravity model discussed in this paper has a more complicated action and the scalaron, which is in fact a complex field and we are focusing on its real part throughout [24], has different masses at different regimes.

Separating this scalaron degree of freedom through appropriate conformal rescaling, the rest of the gravitational action is given by the familiar Einstein term only. The action for a massive scalar field σ , after the conformal rescaling $\sigma \rightarrow e^{-\frac{\kappa\varphi}{\sqrt{6}}} \sigma$ reads

$$\begin{aligned} \mathcal{L}_{\text{scalar}} = & -\frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{\kappa\sigma}{\sqrt{6}} \partial_\mu \sigma \partial^\mu \varphi - \frac{\kappa^2 \sigma^2}{12} \partial_\mu \varphi \partial^\mu \varphi \\ & - \frac{m_\sigma^2}{2} e^{-\frac{2}{\sqrt{6}} \kappa \varphi} \sigma^2. \end{aligned} \quad (66)$$

Similarly, the action for a canonical fermion is given by

$$\mathcal{L}_{\text{fermion}} = -\bar{\psi} \not{D} \psi - e^{-\frac{1}{\sqrt{6}} \kappa \varphi} m_\psi \bar{\psi} \psi, \quad (67)$$

after rescaling $\psi \rightarrow e^{-\frac{\sqrt{6}\kappa\varphi}{4}} \psi$. The spinor covariant derivative is defined, for a U(1) gauge field A_μ , as

$$\not{D} \psi \equiv e^\mu{}_\alpha \gamma^\alpha (\partial_\mu + \Gamma_\mu - ig A_\mu) \psi, \quad (68)$$

$$\Gamma_\mu \equiv \frac{1}{8} [\gamma^\alpha, \gamma^\beta] e^\lambda{}_\alpha \nabla_\mu e_{\lambda\beta}, \quad (69)$$

where $e^\mu{}_\alpha$ is a tetrad (vierbein) field, rescaled as $e^\mu{}_\alpha \rightarrow e^{-\frac{\kappa\varphi}{\sqrt{6}}} e^\mu{}_\alpha$, and Γ_μ is a spin connection for the Dirac spinor [11]. For a U(1) gauge field A_μ and a charged massless complex scalar $\Phi \equiv \chi e^{ig\zeta}/\sqrt{2}$, the transformed action is given by

$$\begin{aligned} \mathcal{L}_{\text{vector}} = & -\frac{1}{4} F^{\mu\nu} F_{\mu\nu} - \frac{1}{2} g^{\mu\nu} \partial_\mu \chi \partial_\nu \chi - \frac{\kappa\chi}{\sqrt{6}} g^{\mu\nu} \partial_\mu \chi \partial_\nu \varphi - \frac{\kappa^2 \chi^2}{12} g^{\mu\nu} \partial_\mu \varphi \partial_\nu \varphi \\ & - \frac{g^2}{2} \chi^2 g^{\mu\nu} V_\mu V_\nu, \end{aligned} \quad (70)$$

where $F_{\mu\nu} \equiv \partial_\mu V_\nu - \partial_\nu V_\mu = \partial_\mu A_\nu - \partial_\nu A_\mu$, a gauge invariant field $V_\mu \equiv A_\mu + \partial_\mu \zeta$ has been defined, and rescaled as $\chi \rightarrow e^{-\frac{\kappa\varphi}{\sqrt{6}}} \chi$. If the gauge symmetry has been broken explicitly or *spontaneously before rescaling*, one could consider the gauge-breaking mass term

$$\mathcal{L}_{\text{g.b.m.}} = -\frac{m_V^2}{2} e^{-\frac{2\kappa\varphi}{\sqrt{6}}} V^2. \quad (71)$$

The scalaron φ can decay into the scalar σ , fermion ψ , and massive vector V_μ via trilinear interactions:

$$\begin{aligned} \mathcal{L}_{3\text{leg}} &= -\frac{1}{\sqrt{6}M_{\text{Pl}}} \sigma \partial^\mu \sigma \partial_\mu \varphi + \frac{m_\sigma^2}{\sqrt{6}M_{\text{Pl}}} \varphi \sigma^2 + \frac{m_\psi^2}{\sqrt{6}M_{\text{Pl}}} \varphi \bar{\psi} \psi + \frac{m_V^2}{\sqrt{6}M_{\text{Pl}}} \varphi V^2 \\ &= \frac{1}{\sqrt{6}M_{\text{Pl}}} \varphi \partial^\mu \sigma \partial_\mu \sigma + \frac{2m_\sigma^2}{\sqrt{6}M_{\text{Pl}}} \varphi \sigma^2 + \frac{m_\psi^2}{\sqrt{6}M_{\text{Pl}}} \varphi \bar{\psi} \psi + \frac{m_V^2}{\sqrt{6}M_{\text{Pl}}} \varphi V^2, \end{aligned} \quad (72)$$

where we have integrated by parts, used $\delta(\sqrt{-g}\mathcal{L}_{\text{scalar}})/\delta\sigma = 0$, and left only linear terms in φ from Eqs. (66), (67) and (71).

The decay rates of the scalaron are then given by [16, 17, 11]

$$\Gamma(\varphi \rightarrow \sigma\sigma) = \frac{\mathcal{N}_\sigma (M_s^2 + 2m_\sigma^2)^2}{192\pi M_{\text{Pl}}^2 M_s} \left(1 - \frac{4m_\sigma^2}{M_s^2}\right)^{1/2}, \quad (73)$$

$$\Gamma(\varphi \rightarrow \bar{\psi}\psi) = \frac{\mathcal{N}_\psi m_\psi^2 M_s}{48\pi M_{\text{Pl}}^2} \left(1 - \frac{4m_\psi^2}{M_s^2}\right)^{3/2}, \quad (74)$$

$$\Gamma(\varphi \rightarrow VV) = \frac{\mathcal{N}_V m_V^4}{48\pi M_{\text{Pl}}^2 M_s} \left(1 - \frac{4m_V^2}{M_s^2}\right)^{1/2} \left(1 - \frac{4m_V^2}{M_s^2} + \frac{12m_V^4}{M_s^4}\right), \quad (75)$$

where \mathcal{N}_σ , \mathcal{N}_ψ and \mathcal{N}_V are the number of modes for each field. Note that we have added gauge-breaking mass term (71) to the vector action (70); otherwise the decay channel to vector modes do not show up classically.

We are interested in the case masses of the decay products are given by a spatially-dependent expectation value of a scalar field ϕ , which is either the Higgs field in the SM or a flat direction field in SUSY, as $m_\sigma = h\phi$, $m_\psi = y\phi$, and $m_V = g\phi$ with h , y , and g representing, scalar, Yukawa, and gauge couplings, respectively. Hereafter we consider the case masses of the decay products are much smaller than the parent particle, so that the phase space suppression factors in the above decay widths are negligible. Then we can classify each mode according to the dependence on ϕ as follows.

$$\begin{aligned} \Gamma(\varphi \rightarrow \sigma\sigma) &= \frac{\mathcal{N}_\sigma M_s^3}{192\pi M_{\text{Pl}}^2} + \frac{\mathcal{N}_\sigma (h\phi)^2 M_s}{48\pi M_{\text{Pl}}^2} + \frac{\mathcal{N}_\sigma (h\phi)^4}{48\pi M_{\text{Pl}}^2 M_s} \\ &\equiv \mathcal{N}_\sigma (\Gamma_{SK} + \Gamma_{SKM} + \Gamma_{SM}), \end{aligned} \quad (76)$$

$$\Gamma(\varphi \rightarrow \bar{\psi}\psi) = \frac{\mathcal{N}_\psi (y\phi)^2 M_s}{48\pi M_{\text{Pl}}^2} \equiv \mathcal{N}_\psi \Gamma_{FM}, \quad (77)$$

$$\Gamma(\varphi \rightarrow VV) = \frac{\mathcal{N}_V (g\phi)^4}{48\pi M_{\text{Pl}}^2 M_s} \equiv \mathcal{N}_V \Gamma_{VM}. \quad (78)$$

Here Γ_{SK} from scalar kinetic term is independent of ϕ and provides the main part of the decay width, while Γ_{SKM} and Γ_{FM} give equal contribution if $h = y$ which is the case in SUSY models. The former is the interference contribution between scalar kinetic and mass terms, and the latter is due to the fermionic mass term. Finally Γ_{SM} and Γ_{VM} , both from their respective mass terms, are higher order in the mass ratio compared with Γ_{SKM} and Γ_{FM} , so that their contribution will be minor, and we do not incorporate them to the subsequent analysis, which turns out to be a good approximation in our case with a large scalaron mass. Note that the equivalence of the above scalaron-decay approach with the Bogoliubov transformation approach to calculate particle production rate have been shown in [12, 14] in the present context and in [31] in a different context.

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References

- [1] A. A. Starobinsky, Phys. Lett. B **91**, 99 (1980).
- [2] A. H. Guth, Phys. Rev. D **23**, 347 (1981); K. Sato, Mon. Not. Roy. Astron. Soc. **195**, 467 (1981); A. D. Linde, Phys. Lett. B **108**, 389 (1982); A. Albrecht and P. J. Steinhardt, Phys. Rev. Lett. **48**, 1220 (1982).
- [3] G. Hinshaw, D. Larson, E. Komatsu, D. N. Spergel, C. L. Bennett, J. Dunkley, M. R. Nolta and M. Halpern *et al.*, arXiv:1212.5226 [astro-ph.CO].
- [4] V. F. Mukhanov and G. V. Chibisov, JETP Lett. **33**, 532 (1981) [Pisma Zh. Eksp. Teor. Fiz. **33**, 549 (1981)].
- [5] E. Komatsu and D. N. Spergel, Phys. Rev. D **63**, 063002 (2001) [astro-ph/0005036].
- [6] K. Maeda, Phys. Rev. D **39**, 3159 (1989).
- [7] A. D. Linde, Phys. Lett. B **129**, 177 (1983).
- [8] G. Dvali, A. Gruzinov and M. Zaldarriaga, Phys. Rev. D **69**, 023505 (2004) [astro-ph/0303591]; L. Kofman, astro-ph/0303614.
- [9] M. Zaldarriaga, Phys. Rev. D **69**, 043508 (2004) [astro-ph/0306006].
- [10] T. Suyama and M. Yamaguchi, Phys. Rev. D **77**, 023505 (2008) [arXiv:0709.2545 [astro-ph]].
- [11] Y. Watanabe, Phys. Rev. D **83**, 043511 (2011) [arXiv:1011.3348 [hep-th]].

- [12] A. Vilenkin, Phys. Rev. D **32**, 2511 (1985).
- [13] M. B. Mijic, M. S. Morris and W. -M. Suen, Phys. Rev. D **34**, 2934 (1986).
- [14] E. V. Arbuzova, A. D. Dolgov and L. Reverberi, JCAP **1202**, 049 (2012) [arXiv:1112.4995 [gr-qc]]; the decay rate of the Ref. [Eq. (54)] agrees with Ref. [16], with $\Gamma_R/2 = \Gamma_{\text{tot}}$.
- [15] H. Motohashi and A. Nishizawa, Phys. Rev. D **86**, 083514 (2012) [arXiv:1204.1472 [astro-ph.CO]].
- [16] D. S. Gorbunov and A. G. Panin, Phys. Lett. B **700**, 157 (2011) [arXiv:1009.2448 [hep-ph]].
- [17] Y. Watanabe and E. Komatsu, Phys. Rev. D **75**, 061301(R) (2007) [gr-qc/0612120]; its formulas in the arXiv version agree with those of [16], with $F_1(v) \equiv f'(v)/\sqrt{1 + 3f'(v)^2/2M_{\text{Pl}}^2} = 2M_{\text{Pl}}/\sqrt{6}$ which corresponds to the limit $f'(\phi)|_{\phi=v}/M_{\text{Pl}} \rightarrow \infty$.
- [18] T. Faulkner, M. Tegmark, E. F. Bunn and Y. Mao, Phys. Rev. D **76**, 063505 (2007) [astro-ph/0612569].
- [19] T. Kunimitsu and J. Yokoyama, Phys. Rev. D **86**, 083541 (2012) [arXiv:1208.2316 [hep-ph]].
- [20] S. V. Ketov, Phys. Lett. B **692**, 272 (2010) [arXiv:1005.3630 [hep-th]].
- [21] S. V. Ketov and A. A. Starobinsky, Phys. Rev. D **83**, 063512 (2011) [arXiv:1011.0240 [hep-th]]; JCAP **1208**, 022 (2012) [arXiv:1203.0805 [hep-th]].
- [22] A. De Simone, H. Perrier and A. Riotto, JCAP **1301**, 037 (2013).
- [23] K. Enqvist, A. Mazumdar, M. Postma and , Phys. Rev. D **67**, 121303 (2003) [astro-ph/0304187].
- [24] S. V. Ketov and S. Tsujikawa, Phys. Rev. D **86**, 023529 (2012) [arXiv:1205.2918 [hep-th]]; it is missing a term $-7f_1 R/30$ in its Eq. (20), which results in not recovering the Einstein-Hilbert term in the high curvature regime. To translate its formulas into correct ones, β_{KT} should be replaced by $10\beta/3$.
- [25] M.Y. Khlopov and A.D. Linde, Phys. Lett. **138B**, 265 (1984); J.R. Ellis, J.E. Kim and D.V. Nanopoulos, Phys. Lett. **145B**, 181 (1984); M. Kawasaki and T. Moroi, Prog. Theor. Phys. **93**, 879 (1995)
- [26] K. Yamamoto, Phys. Lett. B **168**, 341 (1986); D.H. Lyth and E.D. Stewart, Phys. Rev. Lett. **75**, 201 (1995); Phys. Rev. D **53**, 1784 (1996)
- [27] J. Yokoyama, Phys. Rev. D **70**, 103511 (2004) [hep-ph/0406072].
- [28] V. D. Barger, M. S. Berger and P. Ohmann, Phys. Rev. D **47**, 1093 (1993) [hep-ph/9209232].
- [29] N. Seto, S. Kawamura and T. Nakamura, Phys. Rev. Lett. **87**, 221103 (2001) [astro-ph/0108011].

- [30] K. Nakayama, S. Saito, Y. Suwa and J. Yokoyama, Phys. Rev. D **77**, 124001 (2008) [arXiv:0802.2452 [hep-ph]]; JCAP **0806**, 020 (2008) [arXiv:0804.1827 [astro-ph]]
- [31] J. Braden, L. Kofman and N. Barnaby, JCAP **1007**, 016 (2010) [arXiv:1005.2196 [hep-th]].
- [32] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1303.5082 [astro-ph.CO].
- [33] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1303.5084 [astro-ph.CO].
- [34] J. L. Cervantes-Cota and H. Dehnen, Nucl. Phys. B **442**, 391 (1995) [astro-ph/9505069]; F. L. Bezrukov and M. Shaposhnikov, Phys. Lett. B **659**, 703 (2008) [arXiv:0710.3755 [hep-th]]; A. O. Barvinsky, A. Y. Kamenshchik and A. A. Starobinsky, JCAP **0811**, 021 (2008) [arXiv:0809.2104 [hep-ph]]; F. Bezrukov, D. Gorbunov and M. Shaposhnikov, JCAP **0906**, 029 (2009) [arXiv:0812.3622 [hep-ph]]; J. Garcia-Bellido, D. G. Figueroa and J. Rubio, Phys. Rev. D **79**, 063531 (2009) [arXiv:0812.4624 [hep-ph]]; F. L. Bezrukov, A. Magnin and M. Shaposhnikov, Phys. Lett. B **675**, 88 (2009) [arXiv:0812.4950 [hep-ph]]; F. Bezrukov and M. Shaposhnikov, JHEP **0907**, 089 (2009) [arXiv:0904.1537 [hep-ph]]; A. O. Barvinsky, A. Y. Kamenshchik, C. Kiefer, A. A. Starobinsky and C. Steinwachs, JCAP **0912**, 003 (2009) [arXiv:0904.1698 [hep-ph]]; A. O. Barvinsky, A. Y. Kamenshchik, C. Kiefer, A. A. Starobinsky and C. F. Steinwachs, arXiv:0910.1041 [hep-ph]; F. Bezrukov, A. Magnin, M. Shaposhnikov and S. Sibiryakov, arXiv:1008.5157 [hep-ph]; K. Kamada, T. Kobayashi, T. Takahashi, M. Yamaguchi and J. Yokoyama, Phys. Rev. D **86**, 023504 (2012) [arXiv:1203.4059 [hep-ph]]; A. De Simone, M. P. Hertzberg and F. Wilczek, Phys. Lett. B **678**, 1 (2009) [arXiv:0812.4946 [hep-ph]].
- [35] F. L. Bezrukov and D. S. Gorbunov, Phys. Lett. B **713**, 365 (2012) [arXiv:1111.4397 [hep-ph]].